

Pole inflation and attractors in supergravity

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Abstract

In string-derived supergravity theory, Kähler metric of chiral matter fields often has a pole. Such Kähler metric is interesting from the viewpoint of the framework of the pole inflation, where the scalar potential can be stretched out to be flat around the pole for a canonically normalized field and inflation can be realized. However, when Kähler metric has a pole, the scalar potential can also have a pole at the same point in supergravity theory. We study such supergravity models with the pole, and provide numerical analysis of inflationary dynamics and resultant density perturbation. We also examine attractor behavior of our model.

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1 Introduction

Exponentially accelerated expansion of spacetime at the very early Universe, so called inflation, solves various problems in the standard Big Bang cosmology [1]. Nowadays the inflationary cosmological model is the standard paradigm in cosmology. In most of early studies [1–4], it had been thought that the Higgs field in a Grand Unified Theory provides a false vacuum energy for inflationary expansion. Since it had been pointed out that any gauge singlet scalar field with a chaotic initial condition can derive inflationary cosmic expansion if its scalar potential is flat enough [5], a scalar field to drive inflation, called inflaton ϕ , has been thought to be not necessarily a Higgs field related with any gauge theory.

When we refer recent observational results of temperature fluctuation in the Cosmic Microwave Background (CMB) [6,7], so-called Higgs inflation where the Higgs field in the standard model of particle physics with nonminimal coupling to gravity plays a role of inflaton fits the observational data well [8]. The secret of the success of Higgs inflation is the Weyl rescaling from the Jordan frame to the Einstein frame, the exponential function factor for this frame transformation makes the scalar potential in the Einstein frame V_E for the canonical normalized inflaton χ being exponentially flatter as

$$V_E(\chi) \sim \left(1 + e^{-\frac{2\chi}{\sqrt{6}M_P}}\right)^{-2}, \quad (1)$$

with M_P being the reduced Planck mass. An exponential-like potential can particularly well fit the observational data, because for such a model there is an interesting relation

$$n_s \simeq 1 - \frac{2}{N_e}, \quad (2)$$

where n_s and N_e denote the spectral index of its density perturbation and the number of e-folds, respectively. Then, for a sensible number of e-folds $N_e \simeq 50$ to 60, the best fit value of $n_s (\simeq 0.96$ to $0.97)$ can be automatically reproduced. This property can be also seen in the so-called Starobinsky inflation model proposed back in 1980, where inflaton is a scalar degrees of freedom appears through the Weyl rescaling and, as the common result, has an exponential dependent scalar potential [9]. Hence, as is well known, both Higgs inflation and Starobinsky inflation well explain observational results.

The generality of this feature with an exponential function in the scalar potential can be classified and be well understood in the framework of α -attractor model, where the scalar potential is expressed as

$$V_E(\chi) \sim \left(1 - e^{-\sqrt{\frac{2}{3\alpha}} \frac{\chi}{M_P}}\right)^2, \quad (3)$$

with new model parameter α [10–13]. The original Starobinsky inflation model corresponds to $\alpha = 1$. In Ref. [13], the attractor behavior of the parameter α has been shown. Since it was recognized that the essence of above key feature comes from the fact that the kinetic term in the original field or the original frame has a singular pole, this classification of models was extended and those models have been regarded as pole inflation model [14,15]. If a field

has a singularity in the kinetic term, the field is infinitely stretched around the singular point of the kinetic term by canonical normalization. In this case that the scalar potential has no singularity at the kinetic singular point, the scalar potential is also stretched infinitely. Then, we can realize sufficiently flat potential, where the inflation by the canonical field takes place. Thus, the pole inflation is interesting, and some extensions also have been studied (see e.g. Ref. [16, 17]).

Indeed, the Kähler metric of chiral matter fields can often have a pole within the framework of string-derived supergravity theory, which typically leads to the following form of kinetic term,

$$\mathcal{L}_{KE} = K_{\phi\bar{\phi}} |\partial_\mu \phi|^2 = \frac{p}{(1 - |\phi|^2)^2} |\partial_\mu \phi|^2, \quad (4)$$

where $K_{\phi\bar{\phi}}$ is the Kähler metric and p is a real number. This kinetic term has a singular point at $|\phi| = 1$. Here we use the unit with the Planck scale $M_P = 1$. The above Kähler metric can be obtained from the Kähler potential $K = -p \ln(1 - |\phi|^2)$. Thus, the supergravity model with this Kähler potential would be a good candidate for realizing the pole inflation. Starobinsky inflation models in no-scale supergravity have been studied, for example, in Refs. [18–24]. However, the corresponding scalar potential can have a similar pole at the same point, because the F-term scalar potential is given by

$$V = e^K \left[K^{I\bar{J}} (D_I W)(D_{\bar{J}} \bar{W}) - 3 |W|^2 \right], \quad (5)$$

by Kähler potential K and superpotential W , with $D_I W \equiv K_I W + W_I$, and the overall factor e^K is singular at the same point as the kinetic term. That is, the realization of the inflation potential is not trivial within the framework of supergravity models with the singular kinetic term. It is our purpose to study such supergravity inflation models. We study simple models, which are inspired by superstring theory.

In this paper, we provide inflationary potentials consist of F-term scalar potential with a logarithmic Kähler potential, constant superpotential and additional non-vanishing F-term component by another field, and a constant (or sufficiently flat) term in the scalar potential.

This paper is organized as follows. In section 2, we explain our potential. In section 3, we show numerical analysis. We also study attractor behavior. In section 4, we assume another type of additional flat potential in scalar potential. Section 5 is devoted to conclusions and discussions.

2 Inflation potential

In superstring theory, the following type of Kähler potential

$$K = -p \ln(T + \bar{T} - A|\phi|^2), \quad (6)$$

appears in the consequence of compactification of extra dimensions, where T is a Kähler modulus, ϕ is a chiral superfield, and A is a constant coefficient. Here we use the notation that the superfield and its lowest component are denoted by the same letter. We assume the Kähler

modulus T to be stabilized with a sufficiently heavy mass by a certain mechanism, and replace T with its vacuum expectation value $\langle T \rangle$ below such a mass scale. Then, by rescaling ϕ , we obtain the following Kähler potential

$$K = -p \ln(1 - |\phi|^2). \quad (7)$$

The kinetic term of the ϕ is given by Eq. (4), and which has a pole at $|\phi| = 1$. The F-term scalar potential in the framework of supergravity is given by Eq. (5).

In this section, we assume non-vanishing F-term $|F|^2 \neq 0$ of other fields and a constant potential V_C due to a D-term scalar potential and/or explicit supersymmetry breaking effects, and hence we obtain the following scalar potential

$$V = e^K \left[K^{\phi\bar{\phi}} |D_\phi W|^2 - 3 |W|^2 + |F|^2 \right] + V_C. \quad (8)$$

For simplicity, we consider the ϕ -independent superpotential $W = W_0$. In the following analysis, we often use the dimensionless scalar potential

$$\tilde{V} \equiv \frac{V}{|W_0|^2} = \frac{1}{(1 - |\phi|^2)^p} \left(p|\phi|^2 - 3 + \tilde{F}^2 \right) + \tilde{V}_C, \quad (9)$$

with $\tilde{F}^2 = |F|^2/|W_0|^2$ and $\tilde{V}_C = V_C/|W_0|^2$, normalizing the scalar potential (8) by $|W_0|^2$. We also define

$$\tilde{V}_1 = e^K \left(K^{\phi\bar{\phi}} |D_\phi W|^2 \right) / |W_0|^2, \quad \tilde{V}_2 = e^K \left(-3|W|^2 + |F|^2 \right) / |W_0|^2. \quad (10)$$

That is, $\tilde{V} = \tilde{V}_1 + \tilde{V}_2 + \tilde{V}_C$. This potential depends only on the radial component ϕ_r of the $\phi = \phi_r e^{i\theta}$ and is shown in Fig. 1a with special values of \tilde{F}^2 and \tilde{V}_C . The values of \tilde{F}^2 and \tilde{V}_C are chosen in such a way that the minimum lies between $0 < \phi < 1$ and its vacuum energy there vanishes. Since \tilde{V}_1 decreases and \tilde{V}_2 increases divergently toward the pole at $\phi_r = 1$, the potential around the origin tends to be flat compared with the depth of the potential at the stationary point. (Thus, there is an approximate shift symmetry for small ϕ_r region.)

Along the direction of $\theta = 0$, a canonically normalized field χ which has canonical kinetic energy $\mathcal{L}_{KE} = \frac{1}{2}(\partial_\mu \chi)^2$ can be defined by

$$\phi = \tanh \left(\frac{\chi}{\sqrt{2p}} \right). \quad (11)$$

The pole $\phi_r = 1$ corresponds to $\chi = \infty$. In terms of the χ , the potential (9) is expressed by

$$\tilde{V} = \cosh^{2p} \left(\frac{\chi}{\sqrt{2p}} \right) \left[p \tanh^2 \left(\frac{\chi}{\sqrt{2p}} \right) - 3 + \tilde{F}^2 \right] + \tilde{V}_C, \quad (12)$$

and is shown in Fig. 1b. The potential of χ is not flat around the pole, but singular. This is because of the factor e^K in the supergravity scalar potential. This is different from the simplest pole inflation, where $V(\phi)$ is not singular at the pole of the kinetic term [14]. However, different

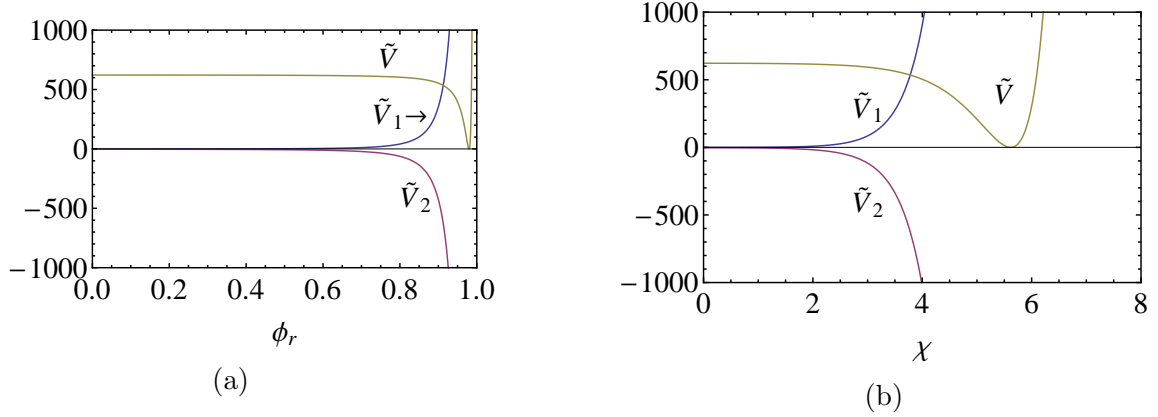


Figure 1: The shape of potential (9) and (12) with $p = 3$, $\tilde{F}^2 = 0.08$, and $\tilde{V}_D = 625$. The pole at $\phi_r = 1$ in Fig. 1a corresponds to the pole at $\chi = \infty$ in Fig. 1b.

singular behaviors between \tilde{V}_1 and \tilde{V}_2 create a minimum and make its depth deeper. Then, the potential for smaller χ region is still flat compared with the depth of the minimum. Around such a flat region, we can realize the inflation.

The stationary condition of the potential $V = |W_0|^2 \tilde{V}$,

$$V_\chi \equiv \frac{dV}{d\chi} = \frac{d\phi}{d\chi} V_\phi = 0, \quad (13)$$

can be expressed by

$$\tilde{V}_\phi = \frac{2p\phi[(p-1)\phi^2 + \tilde{F}^2 - 2]}{(1-\phi^2)^{p+1}} = 0, \quad (14)$$

by using the potential (9). A solution of Eq. (14) is

$$\phi^2 = \frac{2 - \tilde{F}^2}{p - 1}, \quad (15)$$

for $p \neq 1$, and the condition for the stationary point to be present at $0 < \phi < 1$ is

$$3 - p < \tilde{F}^2 < 2. \quad (16)$$

Figure 2 shows the role of \tilde{F}^2 in the potential (12). Smaller \tilde{F}^2 corresponds to deeper vacuum, and the potential becomes sharp around the minimum. Hence, as \tilde{F}^2 becomes smaller, the potential becomes flatter except around the minimum.

3 Numerical analysis

In this section, we study inflation behavior of our potential numerically.

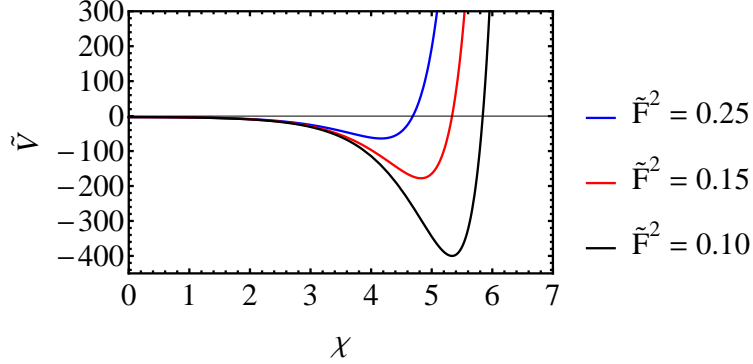


Figure 2: \tilde{F} -dependence of the potential \tilde{V} with $p = 3$ and $\tilde{V}_C = 0$ with respect to the canonical field χ .

3.1 Models with $p = 2, 3$

The shape of potential of inflaton is restricted by observations of the fluctuation of temperature anisotropy in the CMB. Slow-roll parameters are defined by

$$\epsilon = \frac{1}{2} \left(\frac{V_\chi}{V} \right)^2, \quad \eta = \frac{V_{\chi\chi}}{V}, \quad (17)$$

and the power spectrum, the spectral index of the power spectrum of density perturbation, and the tensor-to-scalar ratio are expressed by

$$\mathcal{P}_\zeta = \frac{V}{24\pi^2\epsilon}, \quad n_s = 1 + 2\eta - 6\epsilon, \quad r = 16\epsilon, \quad (18)$$

respectively. According to the BICEP2/Keck Array and Planck data [6, 7], these quantities are measured or bounded as

$$\mathcal{P}_\zeta = (2.20 \pm 1.10) \times 10^{-9}, \quad n_s = 0.9655 \pm 0.0062, \quad r < 0.11. \quad (19)$$

In addition, the number of e-folds

$$N_e \equiv -\ln \left(\frac{a_{\text{end}}}{a} \right) \simeq \int_{\chi_f}^{\chi_{N_e}} \frac{V}{V_\chi} d\chi, \quad (20)$$

is required to be larger than about 50 to solve the flatness and horizon problems. Here χ_{N_e} and χ_f are the field values at the corresponding N_e and the final point of inflation, respectively. In all plots in this paper, we take $\mathcal{P}_\zeta = (2.20 \pm 1.10) \times 10^{-9}$ and $N_e = 60$ for normalization of those quantities.

There are three variables p , \tilde{F} , and \tilde{V}_C in the potential (12). The value of p is quantized such as $p = 1, 2, 3$ in superstring theory. Here, we use $p = 2, 3$ because of the stationary condition (14) in this model. The value of \tilde{F} controls the depth of the potential at the stationary point, which corresponds to the magnitude of flatness of the potential in the region near the origin

too. The value of \tilde{V}_C is adjusted uniquely to vanish the vacuum energy at the minimum. The overall coefficient $|W_0|$ of $V = |W_0|^2 \tilde{V}$ is determined by the amplitude of \mathcal{P}_ζ .

Figure 3 shows (n_s, r) -plots for $p = 2, 3$ with a parameter \tilde{F}^2 . The points (a) and (c) correspond to the lower limit of \tilde{F}^2 in inequality (16), and both of the curves flow to smaller n_s direction as \tilde{F}^2 increases. Table 1 shows variables at each sample point (a)-(e) in Fig. 3. In both cases $p = 2, 3$, there are parameter regions of \tilde{F}^2 consistent with the observations (19). At each point (a)-(e), the superpotential $|W_0|$ takes almost the same scale $10^{-7} M_P^3 \simeq (10^{16} \text{ GeV})^3$, and the constant potential $V_C \simeq (10^{15} \text{ GeV})^4$. The mass of inflaton becomes $m_\chi^2 \simeq (10^{15} \text{ GeV})^2$. The values of \tilde{V}_C , $|W_0|$, and m_χ^2 are analytically calculable by using approximations, and will be derived in section 3.3.

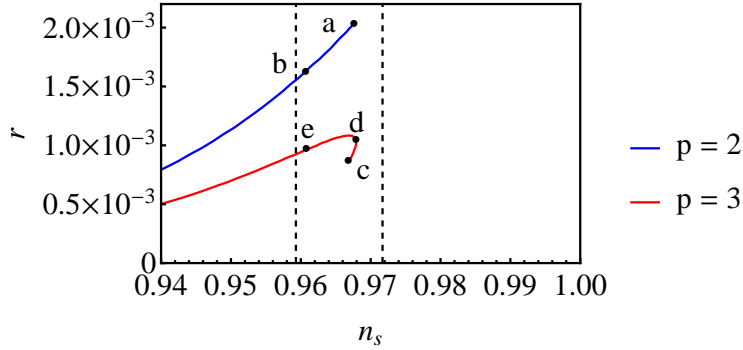


Figure 3: (n_s, r) -plots for $p = 2$ (blue) and $p = 3$ (red). Two vertical dashed lines represent the lower and upper bounds for n_s in the observational data (19). The values of variables at sample points (a)-(e) are showed in Table 1.

Point	p	\tilde{F}^2	\tilde{V}_C	$ W_0 $	n_s	r
a	2	1.001	1.00×10^3	2.59×10^{-7}	0.9676	2.03×10^{-3}
b	2	1.013	7.69×10	8.44×10^{-7}	0.9607	1.63×10^{-3}
c	3	0.005	1.60×10^5	1.33×10^{-8}	0.9670	0.87×10^{-3}
d	3	0.055	1.32×10^4	1.58×10^{-7}	0.9679	1.01×10^{-3}
e	3	0.170	1.38×10^2	4.82×10^{-7}	0.9608	0.96×10^{-3}

Table 1: The values of variables at each point (a)-(e) in Figure 3. The former four variables are input parameters in this model, and the latter two are output quantities.

3.2 Attractor

We showed the numerical analysis for $p = 2, 3$ in Figure 3. In this section, we treat p as not a realistic integer but a continuous parameter in order to examine attractor properties.

Figure 4 shows (n_s, r) -plots by varying p continuously for several values of \tilde{F}^2 . The value of p is bounded by the stationary condition (16). Each curve flows to smaller n_s direction as p increases. It is found that the spectral index n_s becomes realistic values around $3 - p \approx \tilde{F}^2$.

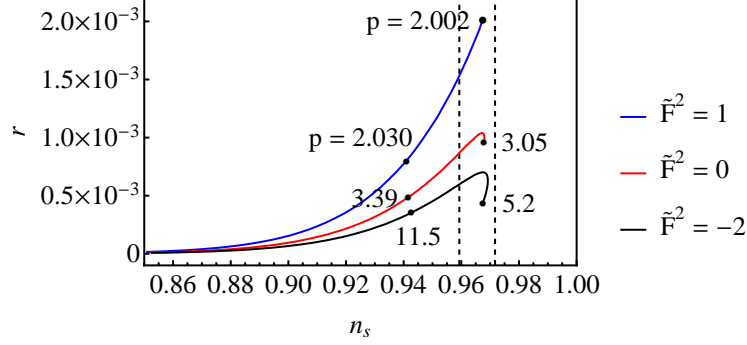


Figure 4: (n_s, r) -plots for different \tilde{F}^2 . The blue, red and black curves represent $\tilde{F}^2 = 1, 0, -2$, respectively. The number written next to each point on curves represents the value of p at the point.

It is convenient to define δ by

$$\delta \equiv \tilde{F}^2 + p - 3. \quad (21)$$

The stationary point (15) and the stationary condition (16) can be rewritten as

$$\phi^2 = 1 - \frac{\delta}{p-1}, \quad 0 < \delta < p-1, \quad (22)$$

respectively. Smaller δ , which makes the stationary point closer to the pole, realizes a deeper vacuum. Figure 5 shows (n_s, r) -plots by varying p continuously for several δ . For a sufficiently large p , the spectral index n_s becomes realistic. These values do not change and approach to $n_s \simeq 0.967$ as p increases. Also, as p increases, r becomes smaller. That is, we find an attractor behavior. The point at $p = 130$ is not a fixed point. For instance, $r = \mathcal{O}(10^{-6})$ can be realized by $p = 1000$ on each curve. These results show that $\tilde{F}^2 \approx 3 - p$ is favorable. In the limit $\tilde{F}^2 \rightarrow 3 - p$ and $\delta \rightarrow 0$, the potential becomes sharp around the minimum. Note that although the limit $\delta \rightarrow 0$ is favorable, point (e) with $\delta = 0.17$ in Table 1 also leads to n_s consistent with the Planck result.

3.3 Analytical calculations

We can examine analytically the attractor behavior shown in section 3.2. At first, we show the attractor behavior. Next, we derive the values of quantities mentioned in the last paragraph in section 3.1.

The potential (12) is rewritten by

$$\tilde{V} = -p \cosh^{2(p-1)} \left(\frac{\chi}{\sqrt{2p}} \right) + \delta \cosh^{2p} \left(\frac{\chi}{\sqrt{2p}} \right) + \tilde{V}_C, \quad (23)$$

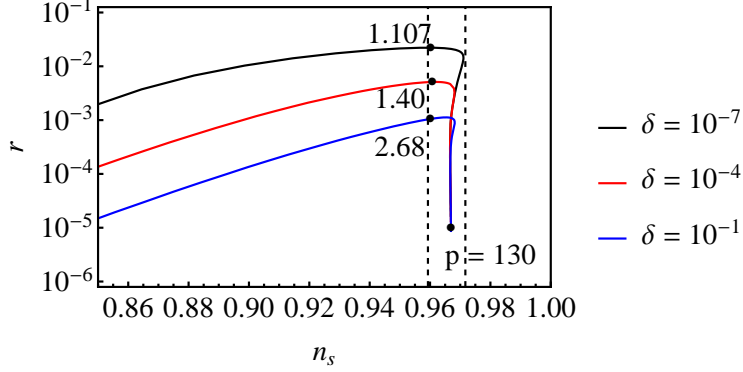


Figure 5: (n_s, r) -plots for different δ . The number written next to each point on curves represents the value of p at the point. At the region near $n_s \simeq 0.967$, three curves almost coincide with each other, and they take almost the same r at $p \gtrsim 100$.

by using δ defined in Eq. (21). We can neglect the factor $\exp(-\chi/\sqrt{2p})$ in $\cosh(\chi/\sqrt{2p})$ in the region related to inflation, and we obtain

$$\begin{aligned}\tilde{V} &= -\frac{1}{2^{2p}} \left(4pe^{-\sqrt{\frac{2}{p}}\chi} - \delta \right) e^{\sqrt{2p}\chi} + \tilde{V}_C, \\ \tilde{V}_\chi &= -\frac{\sqrt{2p}}{2^{2p}} \left[4(p-1)e^{-\sqrt{\frac{2}{p}}\chi} - \delta \right] e^{\sqrt{2p}\chi}, \\ \tilde{V}_{\chi\chi} &= -\frac{1}{2^{2p-1}} \left[4(p-1)^2 e^{-\sqrt{\frac{2}{p}}\chi} - p\delta \right] e^{\sqrt{2p}\chi}.\end{aligned}\tag{24}$$

Being aware of $\tilde{V} \simeq \tilde{V}_C$ during inflation and with Eqs. (24), the number of e-folds is rewritten as

$$N_e = \int_{\chi_f}^{\chi_{N_e}} \frac{V}{V_\chi} d\chi \simeq -\frac{2^{2p-2}\tilde{V}_C}{\sqrt{2p}(p-1)} \int_{\chi_f}^{\chi_{N_e}} d\chi e^{-\sqrt{2p}(1-\frac{1}{p})\chi} \left[1 - \frac{\delta}{4(p-1)} e^{\sqrt{\frac{2}{p}}\chi} \right]^{-1}.\tag{25}$$

In the limit of $\delta \rightarrow 0$ (or $p \rightarrow \infty$), the terms independent of δ in Eq. (25) are dominant and we obtain

$$N_e = \frac{2^{2p-3}\tilde{V}_C}{(p-1)^2} e^{-\sqrt{2p}(1-\frac{1}{p})\chi_{N_e}}.\tag{26}$$

Note that the factor $e^{-\sqrt{2p}(1-\frac{1}{p})\chi_f}$ is negligible because χ_f is sufficiently large. By using (26) with $\delta = 0$, slow-roll parameters (17) are recast as

$$\epsilon|_{\chi=\chi_{N_e}} = \frac{p}{4(p-1)^2 N_e^2}, \quad \eta|_{\chi=\chi_{N_e}} = -\frac{1}{N_e}.\tag{27}$$

Observable quantities (18) are expressed with N_e as

$$n_s = 1 - \frac{2}{N_e} - \frac{3p}{2(p-1)^2 N_e^2}, \quad r = \frac{4p}{(p-1)^2 N_e^2}.\tag{28}$$

With the realistic value of p , e.g. $p = 3$, and $N_e = 60$, we obtain

$$n_s = 0.9664, \quad r = 0.83 \times 10^{-3}. \quad (29)$$

These are consistent with the point (c) in Table 1. For a larger p , instead we obtain

$$n_s \simeq 0.9667 - \frac{1}{p} \times 10^{-3}, \quad r \simeq \frac{1}{p} \times 10^{-3}. \quad (30)$$

The second term in n_s can be neglected for sufficiently large p , and we obtain $n_s \simeq 0.9667$. Eq. (30) is consistent with $(n_s, r) = (0.9669, 1.01 \times 10^{-5})$ at $p = 130$ in Figure 5, and shows attractor behavior.

Next, we estimate the values of \tilde{V}_C , $|W_0|$, and m_χ^2 . The stationary point χ_0 is found from $\tilde{V}_\chi = 0$ in Eqs. (24) as

$$e^{-\sqrt{\frac{2}{p}}\chi_0} = \frac{\delta}{4(p-1)}. \quad (31)$$

Since the constant potential \tilde{V}_C is determined to realize vanishing vacuum energy $\tilde{V}|_{\chi=\chi_0} = 0$, we obtain

$$\tilde{V}_C = \left(\frac{p-1}{\delta} \right)^{p-1}. \quad (32)$$

The size of superpotential $|W_0|$ is determined by the amplitude of the power spectrum \mathcal{P}_ζ . By using Eqs. (19), (27) and (32), we obtain

$$|W_0|^2 = 24\pi^2 \mathcal{P}_\zeta \cdot \frac{p}{4(p-1)^2 N_e^2} \cdot \left(\frac{\delta}{p-1} \right)^{p-1} \simeq \frac{3p\delta^{p-1}}{(p-1)^{p+1}} \times 10^{-11}. \quad (33)$$

The mass of inflation m_χ^2 is obtained by

$$\begin{aligned} m_\chi^2 &= V_{\chi\chi}|_{\chi=\chi_0} = |W_0|^2 \tilde{V}_{\chi\chi}|_{\chi=\chi_0} = \frac{6\pi^2 \mathcal{P}_\zeta p \delta^{p-1}}{N_e^2 (p-1)^{p+1}} \cdot \frac{2(p-1)^p}{\delta^{p-1}} \\ &\simeq \frac{6p}{p-1} \times 10^{-11}. \end{aligned} \quad (34)$$

4 D -term potential

In the previous section, we assume the constant V_C in the potential (8). Since the role of V_C is to uplift the potential to give vanishing vacuum energy, we can consider another uplifting term depending on the field χ . In this section, we show one example, that is the D-term scalar potential written by $V_D = \frac{1}{2}g^2 D^2$ with g being a gauge coupling. For earlier works on such attempts in a different context, see e.g., Refs. [25, 26] for D-term uplifting and Ref. [27] for inflation. When this is independent of ϕ , it corresponds to the constant V_C . Here, we consider the case that the gauge coupling g depends on ϕ as

$$g^{-2} = g_0^{-2} + \beta \ln \left(\frac{\phi}{M_P} \right). \quad (35)$$

The second term depends on ϕ and such a dependence would appear when charged matter fields become massive by its expectation value of ϕ . g_0 and β denote the ϕ -independent part of the gauge coupling and the corresponding beta function coefficient. Also we assume that D itself is independent of ϕ and just a constant. Then, the full scalar potential is written by

$$V = e^K \left[K^{\phi\bar{\phi}} |D_\phi W|^2 - 3 |W|^2 + F^2 \right] + V_D. \quad (36)$$

After redefinition of parameters, we simplify the D-term potential as

$$V_D = \frac{D^2}{1 + \alpha \ln \phi}. \quad (37)$$

Since a sufficiently small α makes this D -term sufficiently flat in order not to disturb inflation, V_D plays a role of those of V_C in the potential (9).

Figure 6 shows (n_s, r) -plots for different fixed α with the parameter \tilde{F}^2 . The point at the center of the figure represents the lower bound for \tilde{F}^2 by inequality (16), i.e. $\delta \rightarrow 0$, and each curve flows from the point to the outside direction of the figure as \tilde{F}^2 increases. Thus, even if we add ϕ -dependent V_D , there is no change in the limit $\delta \rightarrow 0$ from the previous results, and the region with enough small δ is favorable for any value of p and α .

We can discuss more generic form $V(\phi)$. Judging from the above results, we can expect that successful inflationary potential is realized with not only adding the constant term V_C , but also adding sufficiently flat potential around enough small δ .

5 Conclusion

We have studied inflationary dynamics and the resultant density perturbations in supergravity models, whose Kähler metric has a pole and scalar potential also has the pole at the same field value. We found that successful inflation is possible by appropriate choice of parameters. In particular, the parameter region around $\delta \approx 0$ is favorable. For ϕ -dependent $V_D(\phi)$ term in scalar potential instead of constant uplifting, this property is unaffected and the parameter region around $\delta \approx 0$ is still favorable.

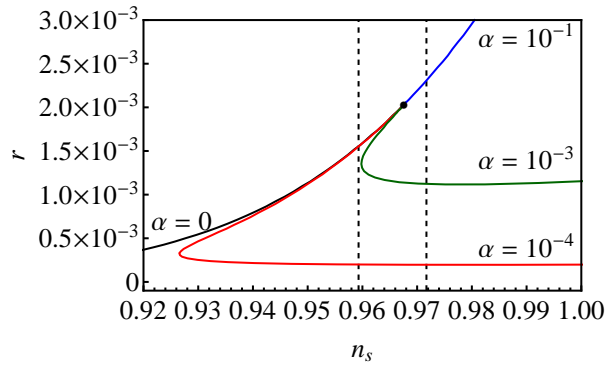
We have also studied the attractor behavior by changing p as a continuous parameter. The parameter region around $\delta \approx 0$ is favorable, again. For a large enough p , as p increases, n_s does not change approaching to $n_s \simeq 0.967$ and only r decreases.

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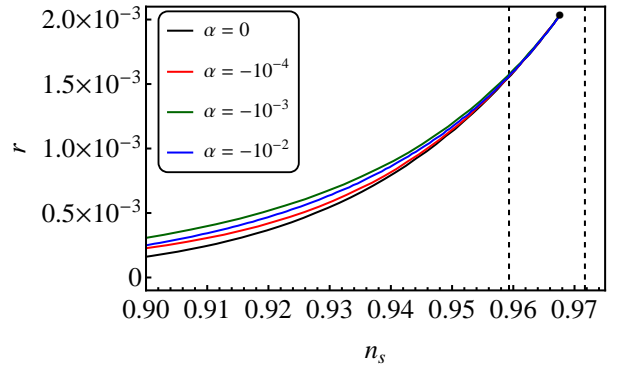
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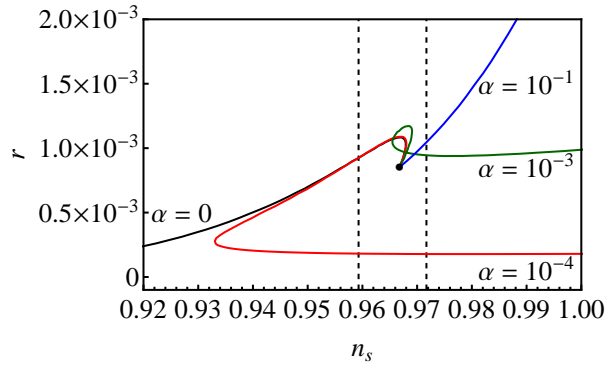
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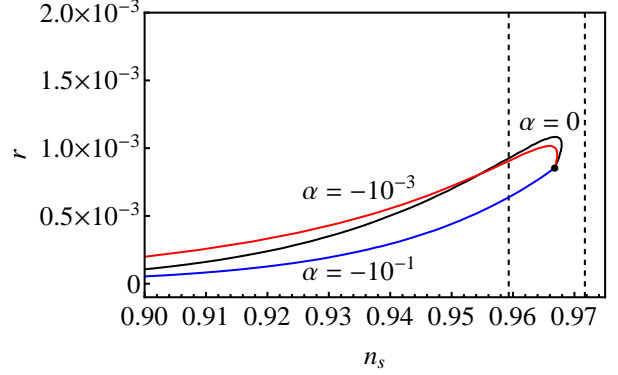
(a) $p = 2, \alpha > 0$



(b) $p = 2, \alpha < 0$



(c) $p = 3, \alpha > 0$



(d) $p = 3, \alpha < 0$

Figure 6: (n_s, r) -plots for $p = 2, 3$ with positive and negative α . The black curve in each figure is the same as Figure 3.

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